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Caviton Collapse of SRS by SBS

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ABSTRACT

Particle simulations are performed to show that the ion waves resulting from Stimulated Brillouin Scatter (SBS) can trigger the collapse of the plasma waves excited by Stimulated Raman Scatter (SRS). We discuss the effect of this collapse mechanism on the hot-electron spectrum produced by the electrostatic waves. A fluid model for the coupling of the ion fluctuations to the plasma waves is formulated and we discuss the necessary condition on the SBS to induce collapse of the plasma waves produced by SRS.

Introduction

The electron plasma waves produced upon Stimulated Raman Scatter from laser-produced plasmas can generate high-energy electrons. As is well-known, these electrons can degrade laser-fusion target gain, and therefore it is crucially important to determine the level and spectrum of these plasma waves. The simplest approach possible is to assume that the waves grow and saturate by wave breaking with the spectrum localized about the linearly most unstable mode. This assumption leads to a simple estimate for the hot-electron temperature, i.e. , that $T_h = 1/2\ v_{ph}^2$, where v_{ph} is the phase velocity of the fastest linearly growing mode of the Raman instability. Such an estimate would no longer be applicable if mode-coupling can become important since the fundamental assumption of localization about the linearly fastest growing mode would be violated. Indeed, recent work on Langmuir turbulence shows that the ion fluctuations caused by the ponderomotive force of the Langmuir waves have a fundamental effect on the wave spectrum.[1] The ion fluctuations provide nucleation points for plasmon wave packets which are driven up by the external source in a manner analogous to a driven resonant oscillator. There is no residual impact of the original linear instability process on the final steady-state spectrum. For the laser-plasma scattering problem an additional source of ion fluctuations can arise directly from SBS. These density waves can directly nucleate electron plasma oscillations. If the ponderomotive force from these driven plasma waves can exceed that originally due to the SBS, the reaction back on the density waves can lead to collapse of the plasma waves. These collapsing plasma waves will then be subject to dissipation through hot-electron generation.

Particle Simulation Results

We have performed simulations in which a large amplitude circularly polarized light wave at frequency ω_o is initialized at t = 0 in a periodic simulation box. Fig.1 summarizes some results of a simulation for the following parameters: $v_o/c = .03$, $v_e/c = .0357$, $M/m = 100$, $\omega_o/\omega_{pe} = 2.5$. In Fig. 1(a) is shown a snapshot of the high-frequency electrostatic field at a relatively early time in the simulation run. This field is predominantly excited at the mode 12, where the scattered wave is at mode 4 for reference and the incident wave is at mode 8. Thus from the matching condition on the ion wave one would expect excitation of the ion wave at mode 16. A result confirmed by Fig. 1(b), a snapshot of the the ion density. In Fig. 1(c) a later-time snapshot of E is shown for comparison. At these later times E has taken on the symmetry of the ion density waves, i.e., the electric field is now phase-locked to the ion waves. In addition to a distinct change in the spatial symmetry of the electric field at this time, there is significant evidence of wave steepening. In Fig. 1(d) we show the electron energy distribution at late time. The spectrum shows evidence of strong electron heating. The same calculation was performed for fixed ions. In contrast to the results recorded in Fig. 1 no wave steepening occurs and there is no evidence of particle heating. Thus one may conclude that the presence of th ion waves from SBS can strongly modify the character of the electron velocity spectrum. Without the SBS coupling the light wave and the two daughter waves of the Raman instability simply enter a cyclical exchance of energy if there is no intrinsic dissipation on the waves. This result follows since the simulation solves the initial value problem . For this fixed ion case If pump energy was continuously supplied to the plasma the waves would have grown to such high intensity that they would eventually result in wave-breaking.

Fluid Model

To establish the conditions necessary for plasma wave collapse, we record here the elements of a fluid model for the ions and electrons and the coupled set of equations for the scattered electromagnetic waves. We first deal with the electromagnetic waves. It is convenient to use the vector potential of the electromagnetic field. An initial or incident electromagnetic wave with the component $A_x^T = 0$ will be considered. Also we will limit our attention to wave propagation in one dimension, i.e., we assume for simplicity that there is only spatial variation in one direction. From Maxwell's equations we write

$$L\,A^T = -\,n_e A^T,$$

where $L = N_o/\,\omega_{pe}{}^2\,(\,\partial^2{}_t - c^2\partial^2{}_x + \omega_{pe}{}^2)$. The total density of electrons is given by $N_o + n_e$, with N_o the background density. We now write $A^T = A^l + A^h$, and $n_e = n_e{}^l + n_e{}^h$. Then on substituting into the equation for A^T, we find that

$$LA^l = -\,n_e{}^l A^l - \langle n_e{}^h A^h \rangle,$$

$$LA^h = -\,n_e{}^h A^h - n_e{}^l A^h - (\,n_e{}^h A^h - \langle\, n_e{}^h A^h \rangle\,).$$

We have not specified what l and h denote at this time. We now rely on the physical assumption that it is possible to separate out a density fluctuation varying on a time scale slow compared with $1/\omega_p$, where $\omega_p{}^2 = 4\pi N_o e^2/m_e$. We denote $n_e{}^l$ as this density fluctuation, and the remainder by $n_e{}^h$. The projection operator $\langle\ \rangle$ projects out all time dependences varying on the same slow time scale. If at this juncture we wish to identify a pump mode we write $A^l = A_o + A^l{}_s$. Now if A_o is assumed large compared to $A^l{}_o$ it is natural to choose the equation which A_o satisfies to have source terms which are second order in the fluctuating quantities. Therefore, we define A_o, and $A^l{}_s$ to satisfy the equations,

$$L\,A_o = -\,n_e{}^l A^l{}_s - \langle n_e{}^h A^h \rangle, \text{ and}$$

$$L\,A^l{}_s = -\,n_e{}^l A_o.$$

Then the coupled set becomes

$$(L + n_e^1)\, A^h = - n_e^h A_o - n_e^h A_s^1 - (n_e^h A^h - \langle n_e^h A^h \rangle)$$

$$L\, A_s^1 = - n_e^1 A_o$$

$$L\, A_o = - n_e^1 A_s^1 - \langle n_e^h A^h \rangle$$

The usual coupled set of equations for the pump and scattered fields follows if we ignore the $n_e^1 A^h$ term on the l.h.s. and the last two terms on the r.h.s. of the first equation of this coupled set.

Alternatively, if we do not separate out A_s^1 , we have

$$(L + n_e^1)\, A^1 = - \langle n_e^h A^h \rangle,$$

$$(L + n_e^1)\, A^h = - n_e^h A^h,$$

if we again ignore the term $(n_e^h A^h - \langle n_e^h A^h \rangle)$ corresponding to higher harmonics. We note that from this approach the contribution due to source n_e^1 enters symetrically in the two equations, and the usual Brillouin instability then is contained in the density term $n_e^1 A^1$. If we set $n_e^1 = n$, $n_e^h = n^h, n^h$ through the ansatz $n^h = -1/8\pi e\ \nabla_x E \exp i\omega_p t + c.c.$, where the amplitude is slowly varying on the time scale $1/\omega_p$, and the total electromagnetic field in terms of the ansatz $A^1 = 1/2 \exp i\omega_o t + c.c.$, and $A^h = 1/2\ A_+ \exp i\omega_+ t + 1/2\ A_- \exp i\omega_- t + c.c.$, with $\omega_\pm = \omega_p \pm \omega_o$, corresponding to light at the pump frequency ω_o, and scattered waves up and down shifted from the pump frequency, we find that the slowly varying amplitudes $A_\pm$, and E are given by the equations,

$$(2i\omega_{\mp}\, \partial_t + (\omega_p^{\,2} - \omega_{\mp}^{\,2}) - c^2\partial_x^{\,2} + \omega_p^{\,2} n/ N_o)\, A_{\mp} = \omega_p^{\,2}/ 8\pi N_o e\, \partial_x E\, A^c, \qquad (1)$$

$$(2i\omega_p\partial_t + \omega_p^{\,2} n/N_o - 3\, v_e^{\,2}\partial_x^{\,2})\, E = -1/2\ \omega_p^{\,2}/c^2 (e/m)\partial_x(A^* A_+ + AA_-), \qquad (2)$$

respectively, and n is the solution of,

$$(\partial_t^{\,2} - c_s^{\,2}\partial_x^{\,2})n = 1/16\pi M\ \partial_x^{\,2}(|E|^2 + \omega_p^{\,2}/ c^2(|A|^2 + |A_+|^2 + |A_-|^2)). \qquad (3)$$

In Eq.(1) the c denotes conjugate for the minus wave only, and In Eq.(3) c_s is the ion sound speed. These equations have been derived from the standard fluid equations for the electron and ion species where electrostatic steepening terms have been neglected. This assumption is generally regarded as justified so long as $n_e/N_o < 1$.[2] We would point out that Eqs. (2) and (3) are analogous to the well-known Zakharov equations,[1] where the source term in Eq. (2) is the result of the beat of A and the scattered waves, and source term in Eq. (3) has ponderomotive force contributions from the electromagnetic wave as well as the usual one due to the electrostatic field. We have not written the equation for A, which would control the pump and the Brillouin scattered wave since in the next section we will consider the case of a fixed pump.

For the sake of orientation with respect to the stimulated scatter instabilities, we note that if we neglect n in Eq. (1), we recover the usual Raman instability. The Brillouin instability, on the other hand, results from the neglect of $A_{\mp}$ and E in the equation for A.

Conditions on Caviton Collapse

Before investigating the conditions for collapse, we first consider the effect of a growing periodic density modulation on SRS. If we ignore the small contribution of A_+, set $A = A_o$, i.e., ignore pump depletion and the contribution to A from the backscattered Brillouin wave, we can write for Eq. (2),

$$(i\partial_t - \partial_x^{\,2})E + nE = -1/\sqrt{3}\, A_o A_- . \qquad (4)$$

For notational convenience we have introduced the normalized units: $t^* = 3/2\ M/m\ \omega_p^{-1}$, $n^* = 4/3\ N_o m/M$, $x^* = 3/2\ \sqrt{M/m}(v_e/\omega_p)$, $E^* = 8\ \sqrt{\pi N_o/3T_e}(m/M)$, and $A^* = \omega_p/c\ E^*$.

As compared with the case of the usual SRS instability, where it is appropriate to employ plane wave states for the spatial variation of E, with nonzero n it is more useful to employ the solutions to the eigenvalue equation,

$$(\lambda + \partial_x^{\,2})\psi - n(x,t)\psi = 0. \qquad (5)$$

For a periodic spatial variation of n, Block's theorem lets us write

$$\psi_k = \exp ikx\ u_k(x), \qquad (6)$$

where u_k has the periodicity of n , and due to the spatially periodic boundary conditions in a box of length L the k's are $k = \pi/\lambda_n(\,2p/N - 1)$, $p = 1, 2, \ldots, N$.[3] Here λ_n is the period length of n and $L = N\lambda_n$. If we use as an example the calculation of the previous section, for which $L = 16\lambda_n$, and λ_n is the wavelength determined by the Brillouin instability, we find that k includes 16 modes, $-7, \ldots, 0, 1, \ldots, 7, 8$ where mode 1 is $2\pi/L$. The various k's are the band labels. Since $u_k(x)$ is periodic with the period of n it can be decomposed into plane waves which are integral multiples of mode 16. Therefore of the 16 possible k modes only mode −4 gives a ψ which has an overlap integral with the r.h.s. of Eq. (5), which is a source at mode 12. We conclude then that ψ and therefore E is composed of the plane wave components at modes $16s - 4$, for integer s.

As a simple approximation for the Bloch state we consider the two lowest energy modes, i.e., the superposition of modes 12 and −4. Then we can write

$$\psi = 1/\sqrt{L}\ (\ \psi_{12} \exp -ik_{12}x + \psi_{-4} \exp -ik_{-4}x\ ), \tag{7}$$

which when substituted in Eq. (5) gives a 2x2 matrix equation for the Bloch states and their eigenvalues. For the eigenvalues we have the solutions

$$\lambda_{\mp} = 1/2\ (\ k_4{}^2 + k_{12}{}^2) \mp (|n(16)|^2 + 1/4\ (\ k_4{}^2 - k_{12}{}^2))^{1/2} \tag{8}$$

To solve for E in Eq. (4), we use the ansatz

$$A_-(x,t) = A_-(t) \exp (\ i\omega_{sc}t - ik_{sc}x\ ), \tag{9a}$$

$$A_o(x,t) = A_o \exp (\ i\omega_{in}t - ik_{in}x), \tag{9b}$$

for the incident and scattered Raman light waves. The relation between ω and k for the light waves can be taken to be the same as when $n = 0$ since the spatial dispersion for light is so large. Also, so long as $\lambda \ll k^2_{De}$, we find k_{sc} as given by the linear frequency-matching condition. Then $\omega_{sc} = k_{12}{}^2$, $k_{sc} = k_4$, $\omega_{in} = 0$, and $k_{in} = k_8$, taking account of the phase convention on the electromagnetic fields and the normalization of units. For E we write

$$E(x,t) = H_+(t) \exp i\int_0^t \lambda_+ dt' \ \ \psi_+(x,t), \tag{10}$$

neglecting the contribution ψ_-, since for early times $\lambda_+ \sim k_{12}{}^2$ and is frequency matched to the source, whereas $\lambda_- \sim k_4{}^2$, and therefore starts out mismatched and the growing n merely increases this mismatch.

Substituting these expressions for E, A_-, A_o into the equations for E and A_- we have the coupled set

$$i\partial_t\ H_+(t) = \alpha(t)\ A_o A_-(t) \exp i\int \Delta\omega\ dt', \tag{11a}$$

$$i\partial_t H_+(t) = \beta(t)\, A_o^* H_+ \exp -i\int \Delta\omega\, dt' , \tag{11b}$$

where $\alpha(t) = ik_{12}(L/3)^{1/2}\psi_{12}^+$, $\beta(t) = 1/\sqrt{L}\ \omega_p/(\omega_o - \omega_p)\ \alpha^*$,
$\psi_{12}^+ = 1/\sqrt{L} \int \exp ik_{12}x\ \psi^+$, and the mismatch, $\Delta\omega = k_{12}^2 - \lambda_+$.
This coupled set can be solved for H_+ and A_- in terms of their initial values to give

$$H_+(t) = H_+(0) \exp \int (\gamma + i/2\ \Delta\omega)\, dt' , \tag{12a}$$

$$A_-(t) = A_-(0) \exp \int (\gamma + i/2\ \Delta\omega)\, dt' , \tag{12b}$$

where the growth rate γ is given as $\gamma = \Gamma^2 - 1/4\ \Delta\omega^2$, and $\Gamma^2 = \alpha\beta\ |A_o|^2 =$

$k_{12}^2/3\ \omega_p/(\omega_o - \omega_p)\ |A_o|^2 |\psi_{12}^+|$.

We can write Γ^2 as

$$\Gamma^2 = \nu_R^2 |\psi_{12}^+| , \tag{13}$$

where ν_R is the usual SRS growth rate. The effect of the density modulation on SRS is now clear. As n grows, $|\psi^+_3|$ initially equal to 1 decreases and γ drops. In addition, growth of the density modulation introduces a nonzero frequency mismatch which has the effect of detuning the instability.

Having now studied the question of the influence of the density wave of SBS on the plasma waves of SRS, we turn to the issue of what is the effect of the electron plasma waves now back on the ion density fluctuation. The direction that the nonlinear evolution of the plasma takes emerges from

The direction that the nonlinear evolution of the plasma takes emerges from this coupling. Drawing from our experience with the Zakharov equations, we can state that the condition for collapse is that pressure due to the ponderomotive force of the electrostatic waves exceeds that due to the beat of the pump and the scattered Brillouin wave. We could derive the necessary condition which would lead to collapse using the above simple model. However we will merely draw some qualitative conclusions here. We see that the fact that a nonzero contribution to this pressure arises from the Raman excited electrostatic waves is due to the density modulation caused by the Brillouin scatter. Thus $|E|^2$ has a term which varies spatially as $\cos k_{12}x$ just as the ponderomotive force from the Brillouin beat. On the other hand the density modulation can actually quench the instability. If this quenching due to detuning occurs at a very early stage the electrostatic waves will not grow to a sufficient level to influence the density and the collapse mechanism will not be triggered. Thus we might expect that if $\nu_R \ll \nu_B$ collapse will not occur. If $\nu_R \gg \nu_B$, however, the electrostatic waves may grown so rapidly that the Brillouin density modulation might have little effect on the nonlinear development of the electrostatic waves (wavebreaking will occur). Now in the weak-coupling limit for the Brillouin instability[2] we find that

$$\nu_R / \nu_B = \sqrt{2}\, T_e^{1/4} g,$$

where g is a slowly varying decreasing function of the background density.(g=1 for 1/4 n_c). Thus for kilovolt plasmas to have wave-breaking dominate, i.e., to satisfy $\nu_R \gg \nu_B$ we must be in the strong-coupling limit for Brillouin and therefore the intensity must be so strong that

$$I \lambda_\mu^2 > 2 \times 10^{13} T_e^{3/2}(\text{kev})\, N_c/N_o \text{ Watts-micron}^2/\text{cm}^2,$$

which is the condition for strong-coupling. For 1/4 micron light, a kilovolt plasma, and a density of .1 critical, this condition corresponds to an intensity greater than 3×10^{15} W/cm^2.

We can conclude from the tentative results presented here that the hot-electron spectrum caused by the Raman instability may require a more complete understanding of the Langmuir turbulence than has been previously believed to be the case from earlier simulation results. In particular, for relatively lower intensities of laser drive the wave-breaking mechanism will have to be replaced by an alternative heating mechanism.

References

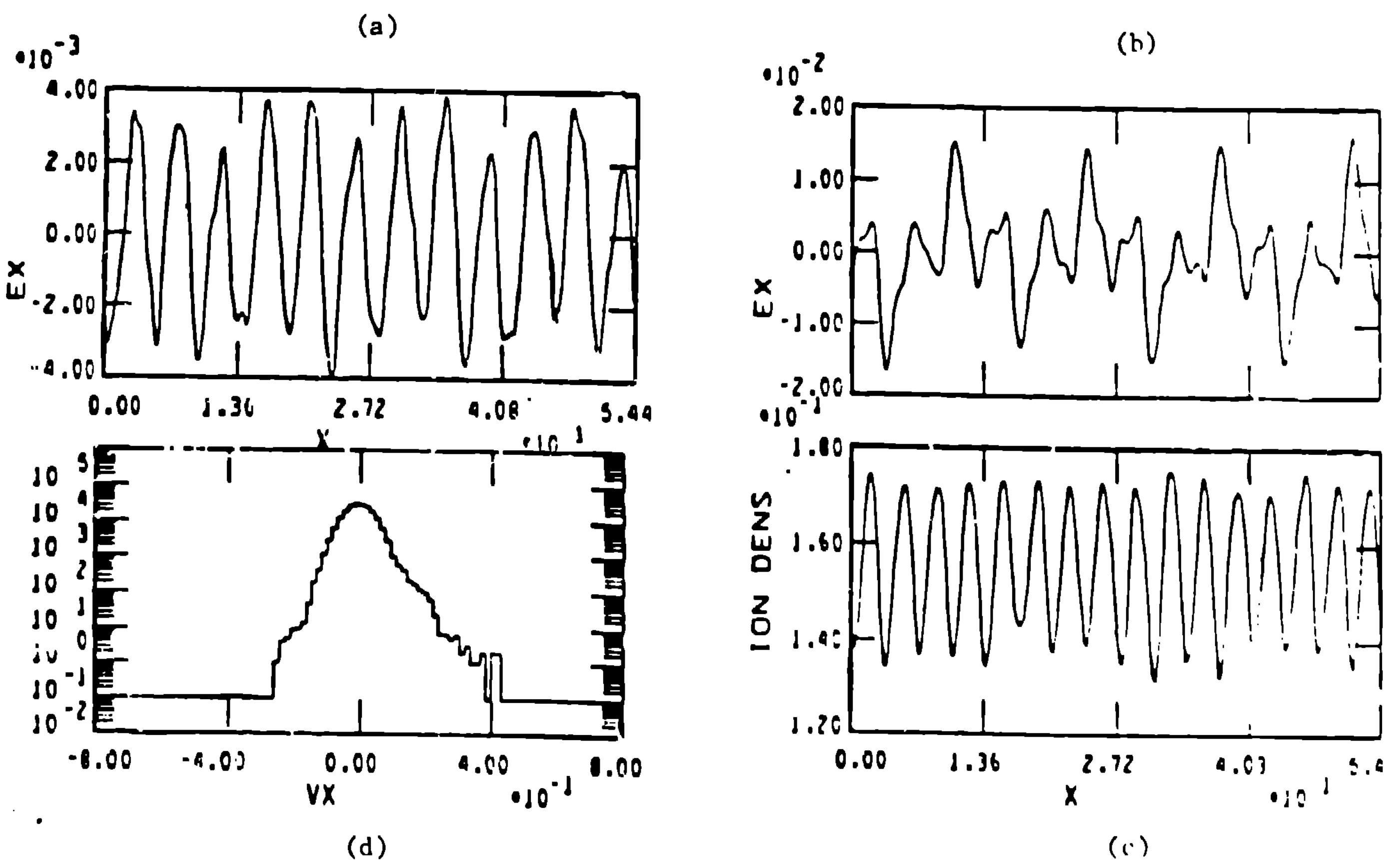
1. G. D. Doolen, D. F. DuBois, and Harvey A. Rose, Phys. Rev. Lett. 54, 804 (1985).
2. D. W. Forslund, J. M. Kindel, and E. L. Lindman, Phys. Fluids, 18, 1002 (1975).
3. R. A. Smith, Wave Mechanics of Crystalline Solids (Chapman and Hall Ltd., London, 1963).

Figure Captions

Figure 1 Periodic particle simulation with following parameters: v_o = .03, v_e = .0357, M/m = 100, Fig. 1(a) is the electrostatic field at time t=596.90. Fig. 1(b) is the electrostatic field at time t =816.81. Fig. 1(c) is ic density at time t = 816.81. Fig. 1(d) is the averaged distribution function. (Units here:

$E_x^* = eE_x/m\omega_o$, $v^* = v_x/c$, N^*/N_c, $x^* = c/\omega_o$.)

Figure 1


(a)
EX
4.00
2.00
0.00
-2.00
-4.00
0.00
2.72
5.44
(b)
EX
2.00
1.00
0.00
-1.00
-2.00
ION DENS
1.80
1.60
1.20
0.00
1.36
2.72
X
VX
-8.00
0.00
4.00
8.00
(d)
(c)